Absence of isentropic expansion in various inflation models

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Dynamics of the inflaton is studied when it interacts with boson and fermion fields and in minimal supersymmetric models. This encompasses multifield inflation models, such as hybrid inflation, and typical reheating models. For much of the parameter space conducive to inflation, the inflaton is found to dissipate adequate radiation to have observational effects on density perturbations and in cases to significantly affect inflaton evolution. Thus, for many inflation models, believed to yield exclusively isentropic inflation, the parameter space now divides into regimes of isentropic and nonisentropic inflation.

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I. INTRODUCTION

The standard picture of inflation [1, 2, 3] has two distinct regimes. In the first phase, the universe undergoes inflationary expansion at constant entropy, so meaning with no radiation production. Thus if inflation started at finite temperature, such as in new inflation scenarios [2], the universe would supercool during the inflation phase, and if inflation started near zero temperature, such as in chaotic inflation scenarios [3], the universe would remain there during the inflation phase. Then in the second phase of inflation, radiation is reintroduced or introduced in a period called reheating. In this letter we show that for a large number of commonly studied inflation models, this picture is flawed. We show that in many models significant radiation production occurs during the inflation period, which has noticeable observational effects on energy density fluctuations as well as in cases completely altering the evolution history of the inflaton.

We examine the model of a scalar inflaton field in interaction with other scalar χ and fermion fields through the basic interaction terms

$$\mathcal{L}_I = -\frac{1}{2} g^2 \phi^2 \chi^2 - g' \phi \bar{\psi}_\chi \psi_\chi - h \chi \bar{\psi}_d \psi_d, \tag{1}$$

where hereafter ϕ is the inflaton field, χ and ψ_χ are additional fields to which the inflaton couples, and ψ_d are light fermions into which the scalar χ-particles can decay $m_\chi > 2m_{\psi_d}$. Aside from the last term, this is the typical Lagrangian used in studies of reheating after inflation [4, 5], or more to the point, if the inflaton is to eventually release its energy into the universe, such interactions are to be expected. Moreover in multifield inflation models, such as hybrid inflation [6], the bosonic interaction plays a basic role in defining the model.

We can qualitatively divide the study of the interactions (1) into two regimes, weak coupling $g, g' \lesssim 10^{-4}$ and moderate to large perturbative coupling $10^{-4} \lesssim g, g' \lesssim 1$. In the weak coupling regime the radiative corrections are too small to affect the flatness of the the inflaton effective potential. In the moderate to large perturbative coupling regime, the radiative corrections from individual terms are significant, so some mechanism, like supersymmetry, must be invoked to cancel these effects and maintain an ultraflat inflaton effective potential. Later in the paper we will show that in minimal SUSY extensions of the typical reheating model, decay channels for the χ or ψ_χ particles are present and the ψ_d field above is simply a representative example. Since in the moderate to strong perturbative regime, reheating models will require SUSY for controlling radiative corrections, Eq. (1) with inclusion of the ψ_d field thus is a toy model representative of the typical reheating model.

Before proceeding, the regime in which the radiation energy density during inflation is below the Hubble scale $\rho_r \lesssim H^4$ will be referred to as isentropic, supercooled, or cold inflation [1, 2, 3, 4, 5, 6] and the other regime $\rho_r > H^4$ will be referred to as nonisentropic or warm inflation [7]. As referred to Eq. (1), this letter will focus on the moderate to large perturbative coupling regime.

In the conventional approach [1, 2, 3, 4, 5, 6], for the effective evolution equation of the inflaton background component, $\varphi \equiv \langle \phi \rangle$, the assumption has been that aside from radiative corrections that modify the φ-effective

potential, $V_{\rm eff}(\varphi)$, this equation is the same as its classical counterpart. However, we have shown in earlier works [8, 9, 10], that in addition to radiative corrections, quantum effects also arise in the φ-effective equation of motion (EOM) in terms of temporally nonlocal terms, which generically lead to dissipative effects. Moreover, although SUSY cancels large quantum effects in the local limit, for the dynamical problem the nonlocal quantum effects can not be cancelled by SUSY. Our previous considerations of inflaton dynamics were limited to nonexpanding spacetime. Here we have extended the calculation to the expanding case (for related earlier works see [11, 12]). We also provide in this work possible physically viable particle physics models, including those based on SUSY, that realize the dissipative dynamics being considered. In this letter we will not delve too much into the (sometimes) evolved technical details of the calculations and the quantum field theory dynamics, for which we refer the interested reader to our previous papers [10] and the most recent one [13]. Here we will focus on the application and relevance of the calculations to inflation and cosmology in general. For this, in Sec. II we start by giving the results for the inflaton effective equation of motion that emerges from a response theory derivation that generalizes to expanding spacetime our previous calculations [10]. Then in Sec. III we present physically viable particle physics models that can realize the typical interactions and decaying mechanism leading to our results, discussing in particular models based on supersymmetry where large quantum corrections, that otherwise could spoils the usual flatness requirements for the inflaton potential, can be kept under control. A minimal SUSY model is explicitly presented and its properties discussed. Finally the conclusions are given in Sec. IV.

II. THE EQUATION OF MOTION AND DISSIPATION IN DE SITTER SPACETIME

Our calculation here is a generalization to expanding spacetime of our previous results [10] and whose details are also reported in [13]. The starting point is the full equation of motion for the inflaton, taken as a homogeneous classical background, $\varphi \equiv \varphi(t)$, which for the couplings shown in Eq. (1) gives

$$\ddot{\varphi}(t) + 3H(t)\dot{\varphi}(t) + \xi {\rm R}(t)\varphi(t) + \frac{dV(\varphi)}{d\varphi} + g^2\varphi(t)\langle\chi^2\rangle + g'\langle\bar{\psi}_\chi\psi_\chi\rangle = 0 \;, \tag{2}$$

where $H(t) = \dot{a}/a$ is the Hubble parameter, ${\rm R}(t)$ the curvature scalar, and ξ the gravitational coupling. $V(\varphi)$ is the tree level potential for the inflaton. In the calculations that follow, we will choose this to be a quartic potential $\lambda\phi^4/4$. To satisfy density perturbation constraints, it is well known the self coupling is tiny, $\lambda \sim 10^{-13}$, which is why in writing Eq. (2) we have neglected the insignificant quantum fluctuation corrections coming from the self-interaction of the inflaton field. Thus for our demonstration of dissipation and radiation production, in the φ effective EOM it is enough to consider the inflaton field as a classical background field interacting with a scalar χ and spinor ψ_χ quantum field. As shown previously [10], the two quantum correction terms in Eq. (2) lead to terms contributing both in the linear regime with respect to the inflaton amplitude from the Yukawa interaction due to the fermionic quantum corrections and in the nonlinear regime from the quadratic coupling of χ to the inflaton.

So far we have not said anything about the role of the ψ_d spinors in Eq. (1) and their contribution in Eq. (2). Their main effect is to dress the propagator for the χ scalar and so they enter nonperturbatively in Eq. (2) through the loops of the χ field. In addition, if we also allow open decay channels of χ into $\psi_d, \bar{\psi}_d$, with decay rate $\Gamma_{\chi\to\psi_d,\bar{\psi}_d}$, the main effect of these spinors is to provide an effective damping in the dressed χ Green's functions. This will in turn reflect back in the φ effective EOM Eq. (2) as an effective dissipation for the inflaton field arising from the $g^2\varphi(t)\langle\chi^2\rangle$ term. If the inflaton is also allowed to decay into the ψ_χ fermions, this can manifest itself in Eq. (2) as a dissipation term due to the inflaton direct decay and will be important for instance in the linear regime for the inflaton amplitude, or more specifically, at the time of reheating when the inflaton may oscillate around the minimum of its effective potential. These latter effects that can come from the fermionic quantum correction in Eq. (2) are not the effects we are interested in here; these effects have been well described in several previous studies of reheating after inflation [4, 5]. Here we are concerned instead with the dissipation identified in Refs. [10, 13], associated with the nonlinear regime for the inflaton field, which can manifest during inflation; these effects mainly arise from the nonlocal quantum corrections due to the χ scalars in Eq. (2). As shown in our previous references, we can then express Eq. (2) in terms of a nonlocal effective EOM, relevant in the nonlinear regime, in the form

$$\begin{aligned}\ddot{\varphi}(t) \;&+\; 3H(t)\dot{\varphi}(t) + \xi{\rm R}(t)\varphi(t) + \frac{dV_{\rm eff}(\varphi(t))}{d\varphi(t)} \\ &+\; 4g^4\varphi(t)\int_{t_0}^{t} dt'\varphi(t')\dot{\varphi}(t')K(t,t') = 0,\end{aligned} \tag{3}$$

where [13]

$$K(t,t') = \int_{t_0}^{t'} dt'' \int \frac{d^3q}{(2\pi)^3} \sin\left[2\int_{t''}^{t} d\tau\omega_{\chi,t}(\tau)\right] \frac{\exp\left[-2\int_{t''}^{t} d\tau\Gamma_{\chi,t}(\tau)\right]}{4\omega_{\chi,t}(t)\omega_{\chi,t}(t'')}, \tag{4}$$

$$\omega_{\chi,t}(\tau) = \left[\mathbf{q}^2 a^2(t)/a^2(\tau) + m_\chi^2 + 2(6\xi - 1)H^2\right]^{1/2}, \tag{5}$$

$m_\chi = g\varphi \gg m_{\psi_d}$, and $\Gamma_{\chi,t}(\tau) \simeq h^2 m_\chi^2/[8\pi\omega_{\chi,t}(\tau)]$ is the decay width of scalars χ of comoving momentum $\mathbf{q}$ into fermions ψ_d. The scale factor has been chosen $a(t) = \exp(Ht)$, since we restrict the analysis to the inflationary, quasi-deSitter regime. $H = \sqrt{8\pi V_{\rm eff}/(3m_{\rm Pl}^2)}$ is the Hubble parameter, with $m_{\rm Pl}$ the Planck mass, and $V_{\rm eff}(\varphi(t))$ is the renormalized effective potential for the inflaton, after the local quantum corrections to $V(\varphi)$ have been taken into account.

The kernel $K(t,t')$, Eq. (4), is obtained by a response theory approximation, which takes advantage of the slow dynamics of the inflaton field so that the dynamics happen effectively in the adiabatic regime. This is equivalent to the closed time path formalism at leading nontrivial order [8, 9], which treats the effect of the field χ on the evolution of $\varphi(t)$. The mode functions for the χ-field that enter in the Green's functions appearing in the response theory expansion of the $\langle\chi^2\rangle$ term in Eq. (2) (see e.g. Ref. [10]) are computed from a WKB approximation, which treats expansion effects and the time variation of the background inflaton field. These mode functions are then used in the loop calculations that determine the effect of the χ quantum corrections to the background inflaton EOM. This WKB approximation is generally valid in the adiabatic regime $d\omega_\chi/d\tau \ll \omega_\chi^2$. In our calculation this adiabatic regime emerges since for the parameter values we will be studying the φ motion is overdamped and since the χ mass is heavy $m_\chi \gg H$. In fact, these two conditions, the slow φ dynamics and the heavy χ mass are enough to assure the validity of a WKB approximation for the χ modes. In addition, the condition $m_\chi \gg H$ implies that curvature effects in the χ field quantum corrections to the background inflaton field are subleading, with dominant term being the Minkowski like corrections. This will be explicitly observed below when we numerically compare several results based on Eq. (3). Note, in the limit $H \to 0$, $a \to$ *constant*, Eq. (4) agrees with the corresponding kernel for nonexpanding spacetime in [10]. The physical origin of the nonlocal (dissipative) term in Eq. (3) is as follows. The evolving background field φ changes the mass of the χ boson which results in the mixing of its positive and negative frequency modes. This in turn leads to coherent production of χ particles which then decohere through decay into the lighter ψ_d-fermions.

When the motion of φ is slow, an adiabatic-Markovian approximation can be applied that converts Eq. (3) to one that is completely local in time, albeit with time derivative terms (for Minkowski spacetime details of this approximation are in [10], while for the expanding spacetime case this is derived in [13]). The Markovian approximation amounts to substituting $t' \to t$ in the arguments of the φ-fields in the nonlocal term in Eq. (3). The Markovian approximation then requires self-consistently that all macroscopic motion is slow on the scale of microscopic motion, thus $\dot\varphi/\varphi, H < \Gamma_\chi$. Moreover when $H < m_\chi$, the kernel $K(t,t')$ is well approximated by the nonexpanding limit $H \to 0$. The validity of all these approximations will be examined below. Combining these approximations, the effective EOM Eq. (3) becomes

$$\ddot\varphi + [3H + \Upsilon(\varphi)]\dot\varphi + \xi {\rm R}\varphi + \frac{dV_{\rm eff}(\varphi)}{d\varphi} = 0, \tag{6}$$

where, by defining $\alpha_\chi = h^2/(8\pi)$,

$$\Upsilon(\varphi) = \frac{\sqrt{2}g^4\alpha_\chi\varphi^2}{64\pi m_\chi\sqrt{1+\alpha_\chi^2}\sqrt{\sqrt{1+\alpha_\chi^2}+1}}. \tag{7}$$

Fig. 1 compares the various approximations for a representative case where $g = h = 0.37$ and the inflaton potential is that for chaotic inflation $V_{\rm eff}(\varphi) = \lambda\varphi^4/4$ with $\lambda = 10^{-13}$ [3]. In Fig 1 evolution has been examined at the final stages of chaotic inflation where we start with $\varphi(t_0 = 0) = m_{\rm Pl}$. The solid line is the exact result based on numerically solving Eq. (3). Plotted alongside this, although almost indiscernible, is the same solution expect using the nonexpanding spacetime kernel (dashed line), obtained by setting $H \to 0$, $a \to$ *constant* in Eq. (4), and the solution based on the adiabatic-Markovian approximation of Eq. (6) (dot-dashed line) for the same parameter set. As seen, the expanding and nonexpanding cases differ by very little and the adiabatic-Markovian approximation is in

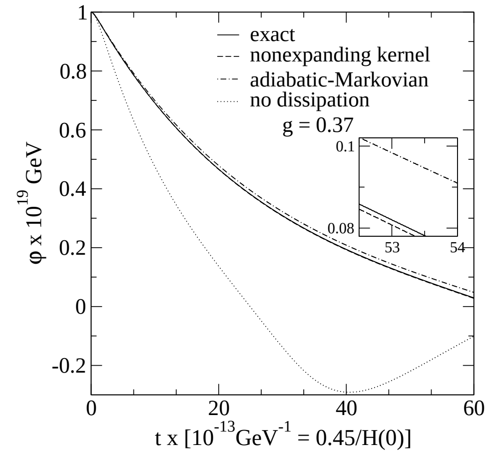

FIG. 1: Evolution of $\varphi(t)$ for $\lambda = 10^{-13}$, $g = h = 0.37$, $\xi = 0$, $\varphi(0) = m_{\rm Pl}$, $\dot{\varphi}(0) = 0$.

good agreement with the exact solution. This confirms simplifying approximations claimed in [9, 10, 11] but up to now had not been numerically verified.

More interestingly, and the first major result of this letter, the dotted line in Fig. 1 is the solution that would be found by the conventional approach in which the nonlocal term in Eq. (3) is ignored. The conventional approach [2, 3, 4, 5, 6] expects the inflaton to start oscillating, which is the precursor to entering various stages of pre/re-heating. However with account for dissipative effects, this never happens for our example in Fig. 1, since the inflaton remains overdamped till the end when it settles at its minima at $\varphi = 0$. Moreover, throughout inflation, and not just near the end, the inflaton dissipates energy, which yields a radiation component of magnitude

$$\rho_r \approx \Upsilon \dot{\varphi}^2/(4H) \,. \tag{8}$$

For the case in Fig. 1, the overdamped regime for the inflaton persists until $g \lesssim 0.35$, below which its evolution at the end of inflation has oscillatory features similar to conventional expectations. However the inflaton is still dissipating radiation at the level Eq. (8) all throughout the inflation period. The radiation energy density produced through this dissipative mechanism is greater than the Hubble scale, $\rho_r > H^4$ and $\Gamma_\chi > H$, if for example $g = h$ and $g > 10^{-2}$ or, as another example, if $h = 0.1$ and $g > 10^{-3}$. If this radiation thermalized, which is expected for $\Gamma_\chi > H$, then $T \approx \rho_r^{1/4} > H$ and it is known that the primordial density perturbations, $\delta\rho/\rho$, produced by the inflaton are altered from their zero temperature value [7, 14, 15, 16]. In particular in the standard relation $\delta\rho/\rho \sim H\delta\varphi/\dot{\varphi}$, for $T < H$, $\delta\varphi^2 \sim H^2$ is the cold inflation result [2, 3, 6], for $T > H$ and $\Upsilon < 3H$, $\delta\varphi^2 \sim HT$ [14, 15], and for $T, \Upsilon/3 > H$, $\delta\varphi^2 \sim \sqrt{H\Upsilon}T$ [16]. Applying these relations to obtain the spectral index n_s for the $\lambda\phi^4/4$ model, for cold inflation the known value is $1 - n_s = 3/N_e$ [17], whereas in contrast for strong dissipative warm inflation ($\Upsilon > 3H$), for the same normalization we find $1 - n_s = 3/(2N_e)$, where N_e is the number of e-folds of inflation. For example for $N_e = 60$ the discrepancy in n_s between the warm and cold inflation cases is about 3%, thus possibly resolvable with current generation high-precision satellite experiments, WMAP and Planck. Moreover cold inflation for any e-folding $N_e > 1$ has the well known η and large φ-amplitude problems, since observationally consistent slow-roll solutions require $m_\phi < H$ and $\varphi > m_{\rm Pl}$ respectively [17]. On the other hand, for warm inflation these problems are nonexistent, since we find for example under the observational constraints $\delta\rho/\rho = 10^{-5}$ and $N_e = 60$, $m_\phi^2 \approx 190H^2 > H^2$ and $\varphi \approx 0.087 m_{\rm Pl} < m_{\rm Pl}$. We thus arrive at our second major result. In multifield inflation models, or in fact any inflation model once reheating interactions are accounted for, there are parameter regimes feasible to inflation, that have never been properly treated since the nonlocal term in Eq. (3) is neglected. Once this term is retained, it is seen that up to fairly weak coupling, adequate radiation is produced during inflation to alter density perturbations. Although this conclusion is based on situations where thermalization is assumed, one can infer the same qualitative outcome once $\rho_r > H^4$, irrespective of the statistical state.

In regimes where Eq. (3) produces significant radiation during inflation, models such as in [5, 6, 18] have to be examined on a case by case basis to determine the effect on density perturbations. Dissipative effects themselves can produce a rich variety of spectra ranging between red and blue [14, 19, 20]. A dramatic example is in [20] which

treats new inflation type symmetry breaking potentials. In cold inflation, such models yield relatively featureless red spectra, whereas with account for dissipative effects, the spectra can be altered between red and blue producing, as an example, the running blue to red spectra suggested by WMAP/2df/Sloan [21].

III. SUSY MODEL IMPLEMENTATIONS AND PARTICLE DECAYING CHANNELS

The above results rely on the ability of the χ-particles to decay, which in our toy model Eq. (1) is fulfilled by the ψ_d-fermions. Based on Eq. (4), this dissipative mechanism is not specific to what the χ particles decay into, but simply requires a nonzero $\Gamma_\chi > H$. Thus the typical possibilities for a scalar field are either decaying into fermions, as in our toy model, or into gauge bosons. Also, in models where χ develops a nonzero vacuum expectation value, such as hybrid inflation [6], the $\phi-\chi$ interaction generates the term $g\langle\chi\rangle\chi\phi^2$, which implies a direct $\chi \to \phi\phi$ decay channel. Independent of inflation issues, particle physics models are adept with interaction configurations similar to Eq. (1). For example, the simplest implementation of the Higgs mechanism in the Standard Model has the background Higgs field coupled to W and Z bosons, thereby generating their masses, and these bosons then are coupled to light fermions through the well known charged and neutral current interactions. In the context of inflation, conventional studies typically only present the inflaton sector itself, since the tacit assumption in these studies has been that interactions are unimportant to the inflation phase. The lesson learned by this work is interactions can play a vital role within the inflation phase, and so care must be taken to understand how the inflaton sector is embedded within the full theory. The few studies that have attempted to do this [18] contain interaction structures similar to Eq. (1).

Whether in our model Eq. (1), reheating models [5], or any other inflation model which has a $\phi - \chi$ coupling in the moderate to large perturbative regime [6], sizable corrections are induced to the φ-effective potential and so SUSY must be invoked. A minimal SUSY extension that incorporates the $\phi - \chi$ coupling would be

$$W = \sqrt{\lambda}\Phi^3 + g\Phi X^2 + fX^3 + mX^2, \tag{9}$$

where $\Phi = \phi + \psi\theta + \theta^2 F$ and $X = \chi + \theta\psi_\chi + \theta^2 F_\chi$ are chiral superfields. In the above model, there would be no additional fermion to associate with ψ_d from our toy model Eq. (1). However the χ-field has a decay channel via a ψ_χ triangle-loop into two light inflaton bosons ϕ. For this case, everything in Eqs. (3)-(5) is unaltered except the decay channel is different with now $\Gamma_\chi \sim (fg^2)^2 m_\chi$. Thus there are additional factors of coupling constants, but in the large perturbative regime, the evolution of φ should be similar to the solid line in Fig. 1. Similarly the ψ_χ-fermion can decay into ϕ and ψ, mediated through a triangle loop comprised of heavy χ and ψ_χ particles.

Other weaker channels also occur. For example, the χ and ψ_χ SUSY partners can have masses differing by $\Delta m \gtrsim m_\phi$, without harming the one-loop cancellations in the φ effective potential. Thus in the model Eq.(9), a χ particle could decay into two fermions, like in our toy model Eq. (1), with one being the fermion partner ψ_χ and the other an inflatino ψ. This channel is phase space suppressed since $m_\chi - m_{\psi_\chi} \sim m_\phi$, which leads to $\Gamma_\chi \sim g^2 m_\phi^2/m_\chi$. Based on Eqs. (6) and (8), this channel leads to $T > H$ for $g \gtrsim 0.3$. Weaker channels such as these need not dominate radiation production but also can serve other particle physics functions such as baryogenesis [7, 23] or dark matter production.

Up to now very little has been said about the role of the ψ_χ-fermion in our toy model. Here its only purpose is to supply a one loop contribution that can suppress the large contributions from the χ sector (then four χ fields are required in Eq. (1)), thus mimic SUSY [9, 10]. For de Sitter spacetime, these results are modified due to ξ-dependent mass corrections and the k_0 integration in this geometry. Both modifications add corrections from the Minkowski spacetime effective potential by terms $\sim O(\ll 1) g^2 H^2 \varphi^2 < m_\phi^2 \varphi^2$ [22] and so can be neglected. Additional analysis of both quantum and thermal corrections for SUSY models like Eq. (9) has been studied independently by the authors of Ref. [24]. They have also concluded that quantum corrections can be kept under control and in addition they find the same holds for thermal corrections.

As briefly discussed below Eq. (8), for warm inflation in the strong dissipative regime, an important result is the avoidance of the η-problem. In cold inflation models, the η-problem is that slow-roll inflation requires $m_\phi < H$ whereas particle physics effects and the coupling to the background metric often prohibit this from happening. In particular in SUSY models, since they imply supergravity, it means higher dimensional operators inherently enter in forms such as $a_n \phi^n / m_p^{n-4}$ or $V\phi^2/m_p^2$, which at either large or small field amplitude can cause the η-problem. However for warm inflation in the strong dissipative regime, slow-roll and density perturbation requirements can be satisfied for $m_\phi > H$ and even $m_\phi \gg H$ [7, 16, 19, 20], thus intrinsically there is no η-problem in any model.

IV. CONCLUSIONS

It is now opportune to give a few remarks on the applications and limitations of our calculation. Eqs. (3) and (4) are valid for any decay channel for the χ-boson, thus any Γ_χ, including the case of no decay channel, $\Gamma_\chi = 0$. If $\Gamma_\chi > H$, then the approximation Eq. (6) also applies. In the regime, $\Gamma_\chi < H$, within the characteristic oscillation period of the free inflaton field $\sim 1/m_\phi$, $K(t,t')$ oscillates considerably, since $m_\chi \gg m_\phi, H$ and without attenuation, since $\Gamma_\chi < H$. Numerically this is a more difficult regime for studying the net effect of the nonlocal term. Another point to note is that Eq. (3) only has the temperature independent contribution to dissipation and so only provides a lower bound. For $\Gamma_\chi > H$ thermalization can in principle be justified and when the temperature is bigger than either the inflaton mass, $T > m_\phi$, or the decay product mass, $T > m_{\psi_d}$ in our toy model, dissipative effects are enhanced [8, 9, 25]. In further work, all these points need to be investigated. Although this work focused on inflaton dissipative effects induced by a χ-boson field, similar considerations can be applied for a ψ_χ fermion. This implies that fermionic and bosonic pre/reheating may be equally affected by our results. In any event, for moderate to large perturbative coupling, SUSY forces both types of interactions to be present side-by-side.

Inflation and reheating phases have been in the past treated as almost detached problems, as so has the quantum-to-classical and the initial condition problems and so on. On the other hand, the warm inflation picture shows, here and elsewhere [7, 19, 20], that inflation can be treated as a single consolidated scenario. As one outcome, various problems, namely η, graceful exit, quantum-to-classical transition, large inflaton amplitude, and aspects of initial conditions [26, 27], can be nonexistent in warm inflation.

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